# Decoherence delays false vacuum decay

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**Abstract**

We show that gravitational interactions between thermal de Sitter modes and a nucleating Coleman-de Luccia bubble lead to efficient decoherence and strongly suppress metastable vacuum decay for bubbles that are small compared to the Hubble radius. The vacuum decay rate including gravity and de Sitter photon interactions has the exponential scaling $\Gamma \sim \Gamma_0^2$, where $\Gamma_0$ is the Coleman-de Luccia decay rate neglecting photon interactions. This strong decoherence effect is a generic consequence of gravitational interactions with light external modes. We argue that efficient decoherence does not occur for the case of Hawking-Moss decay. This observation is consistent with requirements set by Poincaré recurrence in de Sitter space.

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# 1 Introduction

The decay of metastable vacua has been extensively studied and plays a central role in a broad class of cosmological models. The tunneling rate of a single scalar field at a metastable minimum is determined by the bounce solution of the Euclidean equation of motion, as originally demonstrated by Coleman in Ref. [1]. Effects due to coupling to gravity were considered in Ref. [2]. However, in a de Sitter universe there are thermal gravitational modes and realistic cosmologically models have other fields that interact with the tunneling field at least gravitationally. Even though these couplings are Planck suppressed, environmental modes can lead to efficient decoherence, and thus strongly affect the dynamics of a quantum tunneling process.

In this paper we study false vacuum decay, including gravitational couplings to de Sitter modes, considering the specific example of de Sitter photons. Our goal is to determine if the decoherence induced by these interactions is sufficient to significantly suppress the tunneling rate. We find that even though the coupling is Planck suppressed and the wavelength of the de Sitter modes is of order the Hubble radius, decoherence has a significant effect on the vacuum decay rate for vacua that slowly decay via Coleman-de Luccia (CDL) instantons. The decoherence effect can be modeled as a quantum Zeno effect in which the wave function of the tunneling field "collapses" to a classical configuration each time the background leaks information to the environment about whether a bubble exists or not.

Previous works have considered decoherence from modes that are excited by the tunneling field (see e.g. Ref. [3, 4, 5]), taking into account the full master equation that governs the time evolution of the nucleating bubble and all interactions. In this work we restrict ourselves to external modes, so that we can use an S-matrix approach to evaluate the decoherence. This allows us to model the interaction as an ideal partial measurement and greatly reduces the complexity of the problem while keeping a fairly generic form of the interaction. We demonstrate that decoherence due to external modes is far more efficient than decoherence due to modes that are excited by the tunneling field.

The organization of this paper is as follows. In §2 we briefly review how decoherence leads to a delay in the time evolution of a quantum system. Next, in §3 we carefully demonstrate how a field tunneling between two minima in a quantum field theory can be described effectively by a quantum mechanical two-level system using the functional Schrödinger method. We use these results in §4 to determine how decoherence from de Sitter photons influences the bubble nucleation rate. In §5 we remark on the differences between Coleman-de Luccia (CDL) instantons and Hawking-Moss (HM) decay regarding decoherence, and explain how these differences ensure that de Sitter vacua do not survive longer than the recurrence time. We conclude in §6.

## 2 Decoherence and the Quantum Zeno Effect

Let us consider a simple measurement experiment in which a detector is used to determine the state of some two-level system (see e.g. Ref. [6, 7, 8]). Initially, the detector and the system are uncorrelated: $|\psi\rangle = |\psi_{\rm in}\rangle_{\rm det} \otimes |\psi\rangle_{\rm sys}$. Suppose that the interaction Hamiltonian is aligned with the basis $\{|\!\uparrow\rangle_{\rm sys}\,,\ |\!\downarrow\rangle_{\rm sys}\}$. Then after some time we can write

$$|\!\uparrow\rangle_{\rm sys}\,|\psi_{\rm in}\rangle_{\rm det} \quad\rightarrow\quad |\!\uparrow\rangle_{\rm sys}\,|\psi_{\uparrow}\rangle_{\rm det} \tag{1}$$

$$|\!\downarrow\rangle_{\rm sys}\,|\psi_{\rm in}\rangle_{\rm det} \quad\rightarrow\quad |\!\downarrow\rangle_{\rm sys}\,|\psi_{\downarrow}\rangle_{\rm det}\,. \tag{2}$$

Here, we simply relabeled the detector state according to the state it measures. If the two-level system initially is in a coherent superposition $(|\!\uparrow\rangle_{\rm sys} + |\!\downarrow\rangle_{\rm sys})/\sqrt{2}$, we find the reduced density matrix of the measured system by tracing over the detector:

$$\hat{\rho}_{\rm sys} = \frac{1}{2}\begin{pmatrix} 1 & \langle\psi_{\downarrow}|\psi_{\uparrow}\rangle \\ \langle\psi_{\uparrow}|\psi_{\downarrow}\rangle & 1 \end{pmatrix}. \tag{3}$$

Recalling that the off-diagonal entries parametrize the amount of coherence, we immediately see that for $\langle\psi_{\uparrow}|\psi_{\downarrow}\rangle = 0$, all coherence is lost, and the system is reduced to a classical mixture of the two basis states. This matches the intuitive result: once the detector has uniquely determined the state of the system (which corresponds to $|\,\langle\psi_{\uparrow}|\psi_{\downarrow}\rangle\,| = 0$) the wave function "collapses" to one of the eigenstates of the interaction Hamiltonian. To quantify the degree of decoherence that occurs we define the decoherence factor $r$ as

$$r = \langle\psi_{\uparrow}|\psi_{\downarrow}\rangle\,. \tag{4}$$

Note that at no point did we make reference to the size of the detector. It is possible to destroy all coherence of a system if it gets permanently entangled with a single quantum object. In particular, if the detector is entangled with the system and immediately brought out of causal contact we can be certain that the system has lost all coherence. This intuitive observation will turn out to provide a simple mechanism for decoherence in the case of Coleman-de Luccia bubble nucleation.

To see how a quantum Zeno effect arises from interaction with a single quantum object, consider a two-level system that evolves from the state $\Psi_1$ to the state $\Psi_2$ via quantum tunneling. This central system interacts with an environment that is initially uncorrelated. For $t \ll 1/\Gamma$, where $\Gamma$ is the transition rate, this system can be described by the Hamiltonian[1]

$$\hat{H} = \epsilon \hat{\sigma}_z^{\text{sys}} + \Gamma \hat{\sigma}_x^{\text{sys}} + \hat{H}^{\text{env}} + \hat{H}^{\text{int}}. \tag{5}$$

Furthermore, we assume that $\Psi_1$ and $\Psi_2$ are eigenstates of the interaction Hamiltonian, i.e. this is the preferred basis of the environment. We are interested in the decay probability, e.g. the probability for the central system to transition between its two eigenstates after interaction with the environment. Ignoring interactions, one immediately sees that the decay probability for the above Hamiltonian is given by $P_{\text{decay}}(t) = \sin^2(\Gamma t) \approx \Gamma^2 t^2$, where $t \ll 1/\Gamma$ is used in the last approximation.

To be concrete, let the central system initially be in the state $|\Psi_1\rangle$. Thus, the decay probability is given by $P_{\text{decay}} = (1 - \langle \hat{\sigma}_z^{\text{sys}} \rangle)/2$. The time evolution of $\langle \hat{\sigma}_z^{\text{sys}} \rangle$ is then

$$\frac{d\langle \hat{\sigma}_z^{\text{sys}} \rangle}{dt} = i\langle [\hat{H}, \hat{\sigma}_z^{\text{sys}}] \rangle + \left\langle \frac{\partial \hat{\sigma}_z^{\text{sys}}}{\partial t} \right\rangle = 2\Gamma \langle \hat{\sigma}_y^{\text{sys}} \rangle. \tag{6}$$

Considering the intrinsic evolution of the central system, $|\psi(t)\rangle^0_{\text{sys}} = |\Psi_1\rangle - i\Gamma t\, |\Psi_2\rangle + \mathcal{O}(\Gamma^2 t^2)$, we get

$$\frac{d\langle \hat{\sigma}_z^{\text{sys}} \rangle}{dt} \approx -4\Gamma^2 t \ \text{Re}[r(t)]. \tag{7}$$

Thus, for short times the decay probability is given by

$$P_{\text{decay}}(t) = 2\Gamma^2 \int_0^t dt' \ t' \ \text{Re}\left[r(t')\right] + \mathcal{O}(\Gamma^4 t^4). \tag{8}$$

For $r(t) = 1$, the short-time behavior of the isolated system is reproduced. It follows from Eq. (8) that as the decoherence factor approaches zero, the tunneling probability stops increasing. The source of this damping, however, is not immediately obvious. The tunneling rate can be affected when the environment is arranged in such a way that the energy levels of the central system are shifted. Then, the decoherence factor changes by a phase $e^{i\phi(t)}$, and the tunneling probability is affected even though the central system does not get entangled with the environment (e.g. the environment may consist of one-level systems). However,

[1] For simplicity we choose our basis such that $\Psi_1$ and $\Psi_2$ are eigenstates of $\sigma_z$ with opposite eigenvalues.

when an environment is considered that interacts but does not shift the energy levels, the central system leaks information about its state and gets entangled with the environment, such that the absolute value of the decoherence factor decreases. These two processes, which change the survival probability, are complementary.

Note that at no point did we have to make reference to the full master equation for the reduced density matrix that includes the backreaction due to the intrinsic time evolution. This is because we took the preferred basis of the interaction to be aligned with the states between which the central system transitions, i.e. $[\hat{H}^{\rm int}, \hat{\sigma}_z^{\rm sys}] = 0$, and because the interaction lasts only for timescales over which the intrinsic dynamics of the system can be neglected. Let us consider a decoherence factor that decays exponentially with time, say $r \sim e^{-\Gamma_{\rm dec} t}$, which resembles repeated ideal measurements with period $1/\Gamma_{\rm dec}$. In particular, repeated ideal measurements can be described by an S-matrix approach where a detector "scatters" off the system. While these are strong assumptions that do not hold for many scenarios considered in the previous literature (see Ref. [3, 4]), it will turn out that they are satisfied for the interactions considered in this work, namely, gravitational interactions of a true vacuum bubble with massless de Sitter modes.

## 3 Functional Schrödinger Method and Metastable Vacuum Decay

In the previous section we observed how decoherence may lead to suppression of a quantum tunneling process via interactions with the environment[2]. To use the same tools to study bubble nucleation we now carefully match the field theory problem of bubble nucleation to an equivalent quantum mechanics problem.

In the following, we will review the functional Schrödinger method which we will use to derive an effective Hamiltonian that governs the quantum mechanics of the nucleating bubble. The scalar field theory we consider has an $O(4)$-symmetric solution after rotating to Euclidean space. Thus, the full solution can be expressed as a function of one variable, $\lambda$. Once the bubble solution $\phi(\lambda)$ is found, we are interested in how long it will take for the system to tunnel from the metastable vacuum to a field configuration from which the bubble solution can evolve classically. We can treat this as a two-level system in quantum mechanics after obtaining the Hamiltonian for the wave functional $\Psi(\phi)$. Once we obtain the effective Hamiltonian for the intrinsic time evolution of the bubble, we turn to determining the coupling to thermal de Sitter photons. The interaction between the bubble and photons can be treated in an S-matrix approach using the gravitational cross section of a bubble of critical size, which is readily available.

We first carefully examine how the field theory problem can be matched to a quantum mechanical system (we closely follow Ref. [11, 12]). Consider the scalar field theory

$$\mathcal{L} = \frac{1}{2}\partial_\mu\phi\partial^\mu\phi - V(\phi), \tag{9}$$

[2]Possible implications of decoherence in cosmology were considered in e.g. Ref. [9, 10].

where $V(\phi)$ can be any potential. For concreteness we consider the special case of the double well potential

$$V(\phi) = \frac{g}{4}(\phi^2 - c^2)^2 - B(\phi + c). \tag{10}$$

There exists a false vacuum at $\phi = -c$ and a true vacuum at $\phi = c$. The energy difference between the two vacua is approximately $\epsilon \approx 2Bc$. The theory is quantized by demanding that $[\dot{\phi}(\mathbf{x}), \phi(\mathbf{x}')] = -i\hbar\delta^3(\mathbf{x} - \mathbf{x}')$. The resulting functional Hamiltonian is given by

$$H = \int d^3\mathbf{x} \left( -\frac{\hbar^2}{2} \left( \frac{\delta}{\delta\phi(\mathbf{x})} \right)^2 + \frac{(\nabla\phi)^2}{2} + V(\phi) \right). \tag{11}$$

Using the ansatz $\Psi(\phi(\mathbf{x})) = A\exp(-iS(\phi(\mathbf{x}))/\hbar)$ and expanding in powers of $\hbar$ we can write the functional Schrödinger equation at leading order

$$\int d^3\mathbf{x} \left[ \frac{1}{2} \left( \frac{\delta S_0(\phi)}{\delta\phi} \right)^2 + \frac{1}{2}(\nabla\phi)^2 + V(\phi) \right] = E. \tag{12}$$

Note that perpendicular to the most probable escape path (MPEP) the variation of $S_0$ vanishes, while the variation does not vanish along the MPEP. In Ref. [11] it was shown how the WKB method can be used to obtain a solution for $S_0$. We find an effective potential along the MPEP which is parametrized by $\lambda$

$$U(\phi(\mathbf{x}, \lambda)) = \int d^3\mathbf{x} \left( \frac{1}{2}(\nabla\phi(\mathbf{x}, \lambda))^2 + V(\phi(\mathbf{x}, \lambda)) \right). \tag{13}$$

Furthermore, considering that the variation of $S_0$ perpendicular to the escape path vanishes, we find the $O(4)$ symmetric domain wall solution (in the thin-wall approximation)

$$\phi(\mathbf{x}, \lambda) = -c\,\tanh\left( \frac{\mu}{2}(\sqrt{\tau^2 + |\mathbf{x}|^2} - \lambda_c) \right) \approx -c\,\tanh\left( \frac{\mu}{2}(|\mathbf{x}| - \lambda)\frac{\lambda}{\lambda_c} \right) \tag{14}$$

where $\mu = \sqrt{2gc^2}$, $\lambda = \sqrt{\lambda_c^2 - \tau^2}$, and $\lambda_c$ is determined by considering the balance between the domain wall tension $S_1$ and the vacuum energy:

$$S_E = -\frac{\pi^2}{2}\lambda^4 + 2\pi^2\lambda^3 S_1, \tag{15}$$

with the domain wall tension

$$S_1 = \int_{-c}^{c} d\phi\sqrt{2V(\phi)} \approx \sqrt{\frac{g}{2}}\frac{4c^3}{3}. \tag{16}$$

Setting the variation of the total action to zero we find the critical radius of the bubble $\lambda_c = 3S_1/\epsilon$. Any bubble smaller than $\lambda_c$ will decay while any bubble larger than $\lambda_c$ will grow classically. We can use the WKB approximation and the effective potential in Eq. (13) to obtain the solution to the functional Schrödinger equation:

$$\Psi(\phi(\mathbf{x}, \lambda)) = A\exp\left( -\frac{1}{\hbar}\int_0^{\lambda_c} d\lambda\sqrt{2m(\lambda)[U(\lambda) - E]} \right), \tag{17}$$

where the effective mass $m(\lambda)$ is given by

$$m(\lambda) = \int d^3\mathbf{x} \left( \frac{\partial \phi_0(\mathbf{x}, \lambda)}{\partial \lambda} \right)^2 \approx 4\pi S_1 \frac{\lambda^3}{\lambda_c}. \tag{18}$$

We can use Eq. (17) to evaluate the functional for the field configuration of a bubble of critical radius:

$$\Psi(\phi(\mathbf{x}, \lambda_c)) \sim A \exp\left( -\frac{\pi^2}{4\hbar} S_1 \lambda_c^3 \right) . \tag{19}$$

This result deserves some discussion. First, note that while Eq. (19) is just the same exponential scaling as found in Ref. [2], we only solved a time-independent one dimensional quantum mechanics problem[3]. The above argument implies the mapping of the full QFT problem to a quantum tunneling problem: the functional $\Psi$ that has an effective mass $m(\phi(\lambda))$ tunnels through the potential $U(\phi(\lambda))$. This tunneling changes the configuration of the field $\phi$ from the homogeneous false-vacuum solution to a superposition of an expanding bubble and the false vacuum. If one were to measure the state of the field $\phi$ at a time $\tau \ll 1/\Gamma = 2/(\pi^2 S_1 \lambda_c^3)$, the most likely outcome would be the false vacuum solution, while for $\tau \gtrsim 1/\Gamma$ one would most likely observe an expanding bubble.

A possible concern is that interference effects from bubbles of different radii or bubbles at other positions alter the tunneling dynamics. The tunneling rate decreases exponentially with bubble radius. If we are interested in the state of the system at times of order $1/\Gamma(\lambda_c)$, bubbles of smaller radius will have vanished, while bubbles of larger radius have an exponentially suppressed amplitude. Furthermore, for $H \gg \Gamma(\lambda_c)$ only one classically expanding bubble is nucleated per Hubble volume, so that interference effects from other bubbles can be neglected consistently. Of course, this is only true for potentials that do not allow for resonant tunneling, in which case the situation may become more complex (see e.g. Ref. [11]).

We are interested in the time evolution of the tunneling process. As argued above we can treat bubble nucleation as an approximate two-level system that undergoes tunneling from the homogeneous false vacuum configuration to a bubble of critical radius. For times $t \ll 1/\Gamma$ we can define the effective Hamiltonian[4] (see also Ref. [3, 4])

$$\hat{H}_{\text{Bubble}} = \frac{2\pi}{3} R_c^3 \Big( V(\phi_{\text{true}})(1 + \hat{\sigma}_z) + V(\phi_{\text{false}})(\hat{\sigma}_z - 1) \Big) + \Gamma \hat{\sigma}_x \, . \tag{20}$$

In any realistic cosmological model there are more fields than just one isolated scalar. To capture possible effects on tunneling due to environmental degrees of freedom we write the full Hamiltonian in the schematic form

$$\hat{H} = \hat{H}_{\text{Bubble}} + \hat{H}_{\mathcal{E}} + \hat{H}_{\text{int}}, \tag{21}$$

[3]Note that the exponential in Eq. (19) differs from the result for the tunneling rate in Ref. [2] by a factor of two. This is because we calculated the tunneling amplitude rather than the tunneling rate.

[4]We choose our basis such that $\hat{\sigma}_z \Psi_{\text{false}} = +\Psi_{\text{false}}$ and $\hat{\sigma}_z \Psi_{\text{true}} = -\Psi_{\text{true}}$.

where all fields other than $\phi$ are absorbed in the environmental part $\hat{H}_{\mathcal{E}}$. Note that by modeling the bubble as an effective two-level system and neglecting the classical growth after nucleation we underestimate the bubble-photon coupling, and thus obtain a lower bound on the environment induced decoherence.

# 4 Decoherence and False Vacuum Decay

The conclusions of the previous two sections apply for generic bubble-environment interactions that can be modeled by an S-matrix approach, i.e. external modes that interact with the nucleating bubble for a short time during which the intrinsic bubble evolution is negligible. We now turn to a specific environment, consisting of de Sitter photons coupled to gravity, to obtain the decoherence rate and demonstrate the emergence of an efficient quantum Zeno effect.

## 4.1 Particle Interaction

Consider a nucleating bubble $|\mathbf{x}\rangle$ at position $\mathbf{x}$, coupled to an environment of modes $|\chi\rangle_i$ where the interaction is well described by an S-matrix approach (see Ref. [8] for a detailed discussion). Initially, the environment and the bubble are uncorrelated, so the full density matrix factorizes as

$$\hat{\rho}(0) = \hat{\rho}_{\mathrm{B}}(0) \times \hat{\rho}_{\mathcal{E}}(0). \tag{22}$$

We are evaluating the decoherence factor in position space: $r(\mathbf{x}, \mathbf{x}', t)$. This is just the quantity we are interested in, as when coherence over a distance $|\mathbf{x} - \mathbf{x}'| = \lambda_{\mathrm{c}}$ is lost, the MPEP is inaccessible and the bubble nucleation process is highly suppressed. Remember that the decoherence factor is the off-diagonal element of the reduced density matrix, which is given by

$$\hat{\rho}_{\mathrm{B}} = \mathrm{Tr}_{\mathcal{E}}\hat{\rho} = \int d\mathbf{x}d\mathbf{x}' \; \rho_{\mathrm{B}}(\mathbf{x}, \mathbf{x}', 0) \, |\mathbf{x}\rangle \, \langle\mathbf{x}'| \, \langle\chi(\mathbf{x}')|\chi(\mathbf{x})\rangle \, . \tag{23}$$

Assuming no momentum transfer, an isotropic distribution of scattering particles, and a slow intrinsic bubble evolution, the off-diagonal matrix element of the reduced density matrix is determined by (see e.g. Ref. [8])

$$\frac{\partial \rho_{\mathrm{B}}(\mathbf{x}, \mathbf{x}', t)}{\partial t} = -F(\mathbf{x} - \mathbf{x}')\rho_{\mathrm{B}}(\mathbf{x}, \mathbf{x}', t), \tag{24}$$

where

$$F(\mathbf{x} - \mathbf{x}') = \int dq \; \nu(q) v(q) \int \frac{d\mathbf{n}d\mathbf{n}'}{4\pi} \left(1 - e^{iq(\mathbf{n}-\mathbf{n}')\dot{(}\mathbf{x}-\mathbf{x}')}\right) |f(\mathbf{q}, \mathbf{q}')|^2 \, . \tag{25}$$

Here $v(q)$ is the velocity distribution, $\nu(q)$ denotes the momentum density of particles and $|f|^2$ is the scattering amplitude squared. In the long-wavelength limit, the off diagonal component of the density matrix is given by

$$\rho_{\mathrm{B}}(\mathbf{x}, \mathbf{x}', t) = \rho_{\mathrm{B}}(\mathbf{x}, \mathbf{x}', 0)e^{-\Lambda|\mathbf{x}-\mathbf{x}'|^2 t} \, , \tag{26}$$

where

$$\Lambda = \frac{2\pi}{3} \int dq\ \nu(q) v(q) q^2 \left( \int d\cos(\theta)\ \ [1 - \cos(\theta)]\, |f(q,\theta)|^2 \right). \tag{27}$$

Thus, in the long wavelength limit, coherence is lost over a distance $\Delta x$ after times of order $t_{\rm dec} \approx (\Lambda(\Delta x)^2)^{-1}$.

### 4.2 Decoherence Rate of Bubbles via de Sitter Photons

We now use the framework of decoherence developed above to estimate the effects of interactions with thermal photons on bubble nucleation. Note that all assumptions made in Section 2 about the interaction are satisfied for the case of gravitational scattering of photons: the interaction timescale is exponentially small compared to the vacuum decay rate and the preferred basis of the bubble-photon interaction is aligned with the true and false vacuum configuration. At this point it becomes important to check if the decoherence time is small compared to the vacuum decay rate. If this is satisfied we can neglect the intrinsic bubble evolution in the master equation, leading to the simple result for the decay probability found in Eq. (8). It will turn out that decoherence due to external modes are dominant compared to interactions with modes sourced by the tunneling field (see e.g. Ref. [3, 4]).

In order to estimate the decoherence time we evaluate the cross section of gravitational bubble-photon scattering. Let us consider a static, spherically symmetric bubble of true vacuum. In the linear approximation such a configuration leads to the metric ($\eta = \mathrm{diag}(+,-,-,-)$)

$$g_{\mu\nu} = \eta_{\mu\nu} + \kappa h_{\mu\nu}(\mathbf{x}) = \eta_{\mu\nu} - 2\phi(r)(\eta_{\mu\nu} - 2\eta_{\mu 0}\eta_{\nu 0}), \tag{28}$$

where $\kappa^2 = 32\pi G_{\rm N}$ and $\phi$ is the classical potential. Once $\phi(r)$ is fixed we consider the metric to be static. The bubble interacts gravitationally with photons via the action

$$S = -\int d^4x\ \sqrt{g}\frac{F_{\mu\nu}F^{\mu\nu}}{4}, \tag{29}$$

where $F_{\mu\nu} = \partial_\nu A_\mu - \partial_\mu A_\nu$. Expanding $\sqrt{g}$ around flat space gives the vertex for photon-graviton interactions (see Ref. [14]):

$$V_{\mu\nu}(p,p') = \frac{\kappa h^{\lambda\rho}(\mathbf{k})}{2}\left[\eta_{\lambda\rho}p_\nu p'_\mu - \eta_{\mu\nu}\eta_{\lambda\rho}\mathbf{p}.\mathbf{p}' + 2\left(\eta_{\mu\nu}p_\lambda p'_\rho - \eta_{\nu\rho}p_\lambda p'_\mu - \eta_{\mu\lambda}p_\nu p'_\rho + \eta_{\mu\lambda}\eta_{\nu\rho}\mathbf{p}.\mathbf{p}'\right)\right], \tag{30}$$

where $h_{\mu\nu}(\mathbf{k}) = \int d^3\mathbf{x}\ e^{-i\mathbf{k}\cdot\mathbf{x}} h_{\mu\nu}(\mathbf{x})$ is the Fourier transform of $h_{\mu\nu}(\mathbf{x})$. We now turn to evaluating the classical gravitational potential inside a bubble. The most probable size of a non-collapsing bubble is just the critical radius at which the surface tension is balanced by the smaller energy density inside and the gravitational energy. Assuming a bubble of critical radius, the energy in surface tension just cancels the volume energy such that the gravitational potential outside the bubble vanishes. Inside the bubble, the potential is given by

$\phi(r) = \kappa^2 r^2 \epsilon/24$. After Fourier transforming the potential we find the polarization averaged differential cross section to be

$$\begin{aligned} \frac{d\sigma}{d\Omega} &= \frac{1}{(4\pi)^2}\frac{1}{2}\sum_{\text{polarizations}} |\epsilon_r^\mu \epsilon_{r'}^\nu V_{\mu\nu}|^2 \\ &= 64\pi^2 G_{\rm N}^2 |I(R,k)|^2 E^4 (1+\cos(\theta))^2, \end{aligned} \tag{31}$$

where $I(R,k) = \int_0^R dr\, r^2 e^{-ikr} r^2 \epsilon/3$. Note that the cross section at photon momenta $k \ll 1/R$ scales as $\sigma \sim k^4$. Thus, the leading contribution to the decoherence rate is due to modes of wavelengths smaller than the Hubble radius, so that the flat space approximation we used to obtain the scattering amplitude is valid. We can combine the differential cross section in Eq. (31) with Eq. (26) to obtain the decoherence rate. Considering de Sitter radiation at a temperature $T \ll 1/R$ we find

$$\Gamma_{\rm dec} \approx \frac{7\pi \times 2^{16}}{45} G_{\rm N}^2 \epsilon^2 R^{12} T^9. \tag{32}$$

The radius above which a bubble grows classically including gravity is given by $R_{\rm c} = R_0/[1 + (2\pi R_0^2 \epsilon/(3m_{\rm pl}))]$ (see Ref. [2]), where $R_0 = 3S_1/\epsilon$ is the critical radius neglecting gravity. Substituting the critical radius in Eq. (32), the decoherence rate due to thermal de Sitter photons including gravity is given by

$$\Gamma_{\rm dec} \approx \frac{7 \times 3^5 \times 2^7 \sqrt{3}\pi^{10} S_1^{12} m_{\rm pl}^{11} \epsilon^{13/2}}{5\left(6\pi S_1^2 + m_{\rm pl}^2 \epsilon\right)^{12}}, \tag{33}$$

where we used the de Sitter temperature $T = H/(2\pi) = \sqrt{\epsilon/(3m_{\rm pl}^2)}/(2\pi)$. The decoherence rate implied by Eq. (33) is to be compared to the rate of bubble nucleation including gravity, which is given by [2]

$$\Gamma_{\rm CDL} \approx \exp\left(-\frac{27 m_{\rm pl}^4 \pi^2 S_1^4}{2\epsilon(6\pi S_1^2 + m_{\rm pl}^2\epsilon)^2}\right). \tag{34}$$

Note that the approximation in Eq. (34) for the nucleation rate is only valid for cases where the expression inside the exponential is large, such that a polynomial prefactor, corresponding to one-loop corrections, can be neglected. The regime in which the tree level approximation for the nucleation rate is valid is just the same regime where we have $\Gamma_0 \ll \Gamma_{\rm dec}$. This is because the decoherence time is only polynomially small while the nucleation rate is exponentially small. To estimate the rate of bubble nucleation including interactions with de Sitter photons, we can combine Eq. (8), Eq. (33) and Eq. (34) and obtain

$$\Gamma = 2\Gamma_{\rm dec}^{-1}\Gamma_{\rm CDL}^2 \approx \Gamma_{\rm CDL}^2, \tag{35}$$

where we only kept the exponential dependence in the last approximation. This is the main result of this paper. The decoherence induced by interactions with massless external modes leads to an additional factor of 2 in the exponent of the decay rate, indicating a

strong suppression of Coleman-de Luccia bubble nucleation. Furthermore, even though we assumed interactions with de Sitter photons in the above example, the same qualitative features are expected from interactions with de Sitter gravitons. This is because for scattering off a classical gravitational potential, the photon cross section differs only in the angular dependence from the graviton cross section, leading to the same parametric scaling of the decoherence factor (see e.g. Ref. [15]). Note that in deriving the effective nucleation rate we used the thin-wall approximation, which can be written as $S_1^2 \ll V_1 m_{\rm pl}^2$, where $V_1$ is the height of the barrier between the vacua, and we assumed a bubble much smaller than the de Sitter radius.

## 5 Decoherence and de Sitter Recurrence

A possible worry is that any string theoretic description of de Sitter space becomes inconsistent at timescales larger than the recurrence time (see e.g. Ref. [13] and references therein). For a single scalar field $\phi$ the timescale of CDL decay including gravity is given by $t_{\rm decay} \sim e^{S(\phi)+\mathbf{S_0}}$, where $\mathbf{S_0} = -S(\phi_0) = 24\pi^2/V_0 = \log(t_r)$ is the de Sitter entropy and $t_r$ is the recurrence time. If we consider interactions with de Sitter photons, however, we saw in Section 4 that the CDL decay time is changed to about $t_{\rm decay} \sim e^{2(S(\phi)+\mathbf{S_0})}$, which is at risk of exceeding the limits set by Poincaré recurrence. In the following we demonstrate how, despite this apparent inconsistency, the timescale of vacuum decay does not exceed the recurrence time even when interactions with photons and the resulting quantum Zeno effect are included.

There are two possible decay channels through which a false vacuum can decay. For Coleman-de Luccia decay a bubble of true vacuum forms that subsequently grows classically. On the other hand, for Hawking-Moss decay the whole universe tunnels homogeneously out of the false vacuum. In Section 4 we demonstrated how the scattering of external modes provides an efficient mechanism for inducing decoherence. At late times we found a decoherence factor that decreases exponentially with time. This mechanism, which can only occur for CDL tunneling, is very efficient, because after the scattering the detector is out of causal contact with the system, so that coherence cannot be restored. On the other hand, if we consider HM decay in which the whole causal patch tunnels homogeneously, the S-matrix approach is not applicable anymore, as there are no external states. The scenario of continuous system-environment interaction was studied in Ref. [4, 3], where it was found that the decoherence factor decreases polynomially at late times, which is insufficient to induce a strong quantum Zeno effect.

At this point it becomes important to carefully check that a single causal patch can be treated as a closed quantum system that is independent of any physics beyond the horizon. Let us start with the symmetry group of 4-dimensional de Sitter space, $SO(3,1)$. There are three rotations and three boosts. However, only one rotation (spatial rotations) and one boost (time translation) preserve the causal patch. In Ref. [16] it is demonstrated how the

other four symmetries that do not preserve the causal patch are not consistent with assigning a finite amount of entropy to a causal patch in de Sitter space, and thus need to be broken if the holographic principle holds. Thus, from the observer's point of view a causal patch can be treated as an isolated quantum system that does not interact with any degrees of freedom outside the horizon. This indicates that the S-matrix approach to decoherence is not applicable for the case of HM vacuum decay. Hence, we expect the decoherence factor to decrease polynomially with time such that the exponential scaling of the HM vacuum decay rate is not changed by including interactions with other degrees of freedom.

Now that we have shown that HM decay is not significantly affected by decoherence we can reevaluate for what ranges of parameters HM decay dominates over CDL decay, including environmental interactions. Using $\Gamma_{\rm CDL}^{\rm dec} \sim \Gamma_{\rm CDL}^2$ and $\Gamma_{\rm HM}^{\rm dec} \sim \Gamma_{\rm HM}$ as argued above we find (in Planck units)

$$\frac{t_{\rm HM}^{\rm dec}}{t_{\rm CDL}^{\rm dec}} = \exp\left[8\pi^2\left(\frac{16}{S_1^2} - \frac{3}{V_0} - \frac{3}{V_1}\right)\right], \tag{36}$$

which indicates that for $3/V_0 + 3/V_1 > 16/S_1^2$, HM tunneling is the dominant decay channel. The HM decay rate is not changed by decoherence, so any de Sitter vacuum will decay before its lifetime exceeds the limits set by Poincaré recurrence.

## 6 Conclusions

We have demonstrated that the timescale of Coleman-de Luccia decay is highly dependent on external modes to which the tunneling scalar field is coupled. Choosing a generic model of a tunneling scalar field and photons coupled to gravity, we have shown that even de Sitter radiation is sufficient to induce an efficient quantum Zeno effect that suppresses vacuum decay. We exploited the fact that the environmental modes are not sourced by the tunneling field itself, so that we were able to model the bubble-photon interaction using an S-matrix approach. Not only did the use of external modes greatly simplify the problem, it was also a crucial ingredient for obtaining efficient decoherence. While Coleman-de Luccia decay is strongly suppressed, we found that Hawking-Moss decay is not as significantly affected by interactions with the environment. Thus, the lifetime of de Sitter space does not exceed the limits set by the Poincaré recurrence time, even when environmental interactions are included.

The strong suppression of the vacuum decay rate has a broad range of possible implications. In this paper we discussed one specific model of coupling the tunneling field to environmental modes gravitationally. In more realistic cosmological models one expects a far richer pool of fields that couple more strongly to a nucleating new vacuum. We suggest that a far greater suppression of the vacuum decay rate is achievable in such scenarios, e.g. by considering couplings to dark matter or CMB photons. It would be interesting to characterize what the constraints on the stability of de Sitter vacuum are when these decoherence effects are included. In particular, one might expect an effective decay rate that is increasing

with time as the universe gets more and more dilute and decoherence loses efficiency.

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## References

[1] S. R. Coleman, "The Fate of the False Vacuum. 1. Semiclassical Theory," Phys. Rev. D **15**, 2929 (1977) [Erratum-ibid. D **16**, 1248 (1977)].

[2] S. R. Coleman and F. De Luccia, "Gravitational Effects on and of Vacuum Decay," Phys. Rev. D **21**, 3305 (1980).

[3] C. Kiefer, F. Queisser and A. A. Starobinsky, "Cosmological Constant from Decoherence," Class. Quant. Grav. **28**, 125022 (2011) [arXiv:1010.5331 [astro-ph.CO]].

[4] F. Queisser, "Influence of nontrivial backgrounds and decoherence on vacuum decay," arXiv:1004.2921 [gr-qc].

[5] S. Matsumoto, K. I. Nagao, M. Nakamura and M. Senami, "Fate of the False Vacuum Revisited," arXiv:1009.1927 [hep-ph].

[6] W. H. Zurek, "Decoherence, einselection, and the quantum origins of the classical," Rev. Mod. Phys. **75**, 715 (2003)

[7] M. Schlosshauer, "Decoherence, the measurement problem, and interpretations of quantum mechanics," Rev. Mod. Phys. **76**, 1267 (2004) [quant-ph/0312059].

[8] M. Schlosshauer, "Decoherence and the Quantum-to-Classical Transition" (Springer, Berlin/Heidelberg, 2007)

[9] S. Winitzki, "Age-dependent decay in the landscape," Phys. Rev. D **77**, 063508 (2008) [arXiv:0712.2192 [hep-th]].

[10] L. M. Krauss, J. Dent and G. D. Starkman, "Late Time Decay of the False Vacuum, Measurement, and Quantum Cosmology," Int. J. Mod. Phys. D **17**, 2501 (2009) [arXiv:0810.4574 [gr-qc]].

[11] S. -H. H. Tye and D. Wohns, "Resonant Tunneling in Scalar Quantum Field Theory," arXiv:0910.1088 [hep-th].

[12] K. M. Bitar and S. -J. Chang, "Vacuum Tunneling And Fluctuations Around A Most Probable Escape Path," Phys. Rev. D **18**, 435 (1978).

[13] S. Kachru, R. Kallosh, A. D. Linde and S. P. Trivedi, "De Sitter vacua in string theory," Phys. Rev. D **68**, 046005 (2003) [hep-th/0301240].

[14] A. Accioly and R. Paszko, "Photon mass and gravitational deflection," Phys. Rev. D **69**, 107501 (2004).

[15] P. C. Peters, "Differential cross sections for weak-field gravitational scattering," Phys. Rev. D **13**, 775 (2003)

[16] N. Goheer, M. Kleban and L. Susskind, "The Trouble with de Sitter space," JHEP **0307**, 056 (2003) [hep-th/0212209].